

# Neutrino Masses, Baryogenesis and Bilinear R-parity Violation

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## Abstract

We consider the impact of cosmological  $B-L$  constraints on supersymmetric standard models with bilinear breaking of R-parity. In order to avoid erasing any primordial baryon or lepton asymmetry above the electroweak scale,  $B-L$  violation for at least one generation should be sufficiently small. Working in the context of models with non-universal soft supersymmetry breaking masses, we show how the above cosmological constraint can be satisfied while simultaneously providing a neutrino mass matrix required by current data.

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In recent years the increasingly strong evidence for neutrino oscillations from various experiments [1] has led to the active study of R-parity violating extensions of the minimal supersymmetric standard model (MSSM) [2]. Such models maintain the particle spectrum of the MSSM but contain renormalizable lepton flavour violating couplings. The observed neutrino oscillations and mass differences [3] can be accommodated with such couplings [4, 5], and so these models provide a conceivable alternative to seesaw mechanisms [6] of neutrino mass generation. In contrast to the R-parity conserving MSSM, the lightest supersymmetric particle is unstable and decays in the detector with branching ratios which are correlated with the neutrino mixing [7]. This provides a robust, experimentally accessible test of the model at the Large Hadron Collider and/or a  $e^+e^-$  Linear Collider [8]. Analogous confirmatory signatures are less readily found for the elegant seesaw mechanism [6]. Bilinear R-parity violation (BRpV) is the minimal extension of the MSSM with R-parity violating terms [4, 9, 10, 11]. The minimal supergravity version of BRpV [12] (i.e. imposing universal soft supersymmetry breaking masses at an ultraviolet scale) can easily accommodate the atmospheric neutrino oscillation data. However, in order to provide the currently favoured large mixing angle solution for the solar neutrino anomaly, this universality condition must be relaxed [13, 14, 15]. Another option for obtaining a realistic neutrino mass matrix is to allow both bilinear and trilinear couplings while keeping the universality condition of the soft supersymmetry breaking masses. The minimal model of trilinear R-parity violation (TRpV) assumes the dominance of the third generation trilinear couplings and thus contains five free parameters of lepton number violation to fit all the neutrino data successfully [5].

The theoretical background on massive Majorana neutrinos and lepton violating mixing matrices describing neutrino oscillations can be found in [17]. The atmospheric neutrino data is explained by oscillations  $\nu_\mu \longleftrightarrow \nu_\tau$ , and a global analysis gives the following  $3\sigma$  ranges [18]

$$\begin{aligned} 0.3 &\leq \sin^2 \theta_{\text{atm}} \leq 0.7 \\ 1.2 \times 10^{-3} \text{ eV}^2 &\leq \Delta m_{\text{atm}}^2 \leq 4.8 \times 10^{-3} \text{ eV}^2 \end{aligned} \quad (1)$$

with maximal mixing  $\sin^2 \theta_{\text{atm}} = 0.5$  and  $\Delta m_{\text{atm}}^2 = 2.5 \times 10^{-3} \text{ eV}^2$  as the best fit point. Similarly, the solar neutrino data is explained by  $\nu_e$  oscillation into a mixture of  $\nu_\mu$  and  $\nu_\tau$ . Global analysis suggests a large mixing angle, although not maximal, and a much smaller mass squared difference. The allowed region for  $\Delta m_{\text{sol}}^2$  previous to KAMLAND results [1] is

now split into two sub-regions. At  $3\sigma$  we have [18, 19]

$$\begin{aligned} 0.29 &\leq \tan^2 \theta_{\text{sol}} \leq 0.86 \\ 5.1 \times 10^{-5} \text{ eV}^2 &\leq \Delta m_{\text{sol}}^2 \leq 9.7 \times 10^{-5} \text{ eV}^2 \\ 1.2 \times 10^{-4} \text{ eV}^2 &\leq \Delta m_{\text{sol}}^2 \leq 1.9 \times 10^{-4} \text{ eV}^2 \end{aligned} \quad (2)$$

with  $\tan^2 \theta_{\text{sol}} = 0.46$  and  $\Delta m_{\text{sol}}^2 = 6.9 \times 10^{-5} \text{ eV}^2$  as the best fit point.

In connection with neutrino physics, there appears an important cosmological consideration. As is well known, the seesaw mechanism provides a natural way to generate the baryon asymmetry of the universe through the out-of-equilibrium decay of a heavy right-handed neutrino [20]. Being a new physics model just around TeV scale, the R-parity violating MSSM can hardly accommodate such a mechanism of baryogenesis. However, in the supersymmetric model, the so-called Affleck-Dine mechanism can successfully work to generate the required amount of the baryon asymmetry in the flat direction along, e.g.,  $LH_u$  [21]. It is notable that such a property is unaltered even with the presence of R-parity violating terms which must be very small to generate tiny neutrino masses.

It is known that lepton number violating couplings have important consequences for baryogenesis since together with  $B + L$  violating sphaleron processes they are capable of erasing any pre-existing baryon/lepton asymmetry in the universe [22, 23, 24]. The purpose of this paper is to explicitly check if such cosmological constraints on the lepton violating couplings can be satisfied in BRpV while simultaneously accommodating the form of the neutrino mass matrix indicated by the atmospheric, solar and reactor neutrino experiments. A previous analysis [25] derived the cosmological bounds for BRpV but their effect on the neutrino mass matrix was not covered. Given the wealth of new data which has become available since [25] appeared, we develop their analysis and apply the bounds to the currently favoured bimaximal mixing form of the neutrino mass matrix.

We note that our investigation is not relevant if electroweak baryogenesis [26] is operative, in which case the produced baryon asymmetry cannot be erased solely by R-parity violating processes. For our purposes we assume that a  $B - L$  asymmetry was generated primordially by some means at a high energy scale, and our intention is its preservation at all energies down to the electroweak scale when the sphalerons finally fall out of equilibrium.

We briefly summarize the mechanism of neutrino mass and mixing generation by R-parity violating couplings, both bilinear and trilinear. The R-parity violating MSSM predicts a hi-

erarchical neutrino mass spectrum. The atmospheric mass scale corresponds approximately to the heaviest neutrino mass,  $m_3$ , and it is generated at tree level via a low energy see-saw mechanism due to the mixing of the neutrinos with the neutralinos. On the other hand, the solar mass scale, corresponds approximately to the second heaviest neutrino,  $m_2$ , and is generated at the one loop level. The atmospheric neutrino mixing is also predicted by tree level physics, and depends in a simple way on sneutrino vacuum expectation values expressed in the basis where the bilinear parameters are removed from the superpotential. On the other hand, the solar neutrino mixing angle is again predicted by one-loop physics which is mainly determined either by the trilinear couplings in the superpotential or by the bilinear parameters in the scalar potential.

Let us remark, however, that we cannot exclude the possibility of the loop mass dominating over the tree mass, which may have an interesting implication to baryogenesis as will be discussed later.

The well-known baryogenesis constraint [22, 23] can be easily applied to the TRpV model with the universality to exclude this possibility. To see this, let us consider the following trilinear R-parity violating couplings in the superpotential;

$$W = \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c \quad (3)$$

which generates a neutrino mass at one-loop level as follows;

$$M_{ij}^{loop} = 3 \frac{\lambda'_{i33} \lambda'_{j33}}{8\pi^2} \frac{m_b^2 (A_b + \mu \tan \beta)}{m_{b_1}^2 - m_{b_2}^2} \ln \frac{m_{\tilde{b}_1}^2}{m_{\tilde{b}_2}^2} + \frac{\lambda_{i33} \lambda_{j33}}{8\pi^2} \frac{m_\tau^2 (A_\tau + \mu \tan \beta)}{m_{\tilde{\tau}_1}^2 - m_{\tilde{\tau}_2}^2} \ln \frac{m_{\tilde{\tau}_1}^2}{m_{\tilde{\tau}_2}^2}. \quad (4)$$

Note that we have picked up  $\lambda'_{i33}$  and  $\lambda_{i33}$  which give the largest contribution to the neutrino masses when all the trilinear couplings are of similar magnitude. Then, requiring the above one-loop mass (4) gives rise to the solar neutrino mass scale,  $m_2 \approx \sqrt{\Delta m_{sol}^2} \approx 8 \times 10^{-3}$  eV, we obtain

$$\lambda'_{i33}, \lambda_{i33}/\sqrt{3} \approx 5 \times 10^{-5} \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^{1/2} \left( \frac{m_2}{8 \text{ meV}} \right)^{1/2} \quad (5)$$

taking  $\tilde{m} = A_b + \mu \tan \beta = m_{\tilde{b}_{1,2}} = A_\tau + \mu \tan \beta = m_{\tilde{\tau}_{1,2}}$ . Now, the problem is that such a large coupling makes lepton number violating interactions very active when the  $B + L$  violating sphaleron interaction is also in thermal equilibrium, so together they erase the baryon asymmetry before the electroweak phase transition. Indeed, the interaction in Eq. (3)

gives the decay width for lepton number violating one-to-two body decays,

$$\Gamma_{12} = \frac{\pi \lambda_{i33}^{(\prime)2}}{192\zeta(3)} \frac{\tilde{m}^2}{T} \quad (6)$$

assuming  $T \gg \tilde{m}$ . The out-of-equilibrium condition,  $\Gamma_{12} < H = 1.66\sqrt{g_{eff}}T^2/m_{Pl}$ , gives

$$\lambda'_{i33}, \lambda_{i33} < 2 \times 10^{-7} \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^{1/2} \quad (7)$$

for  $g_{eff} = 915/4$ . This is for  $T \gg \tilde{m}$ . An improved result which does not make this assumption was presented in [24] and shows that the  $T/\tilde{m}$  dependence of Eq. (6) is very mild. One sees a big contradiction between (5) and (7). As indicated in Eq. (5), one needs the trilinear couplings of  $\lambda'_{233,333} \sim \lambda_{133,233} \sim \mathcal{O}(10^{-5})$  to accommodate the required bi-large mixing of the atmospheric and solar neutrino oscillation [5]. Thus, the baryogenesis constraint rules out a purely TRpV explanation of the observed neutrino data.

The situation may be different if the neutrino masses are generated purely by bilinear R-parity violating couplings with non-universal soft masses, in which case the non-universality can give much freedom. Forbidding the lepton number violating trilinear couplings in the superpotential in Eq. (3), the BRpV model allows the following dimension-two terms in the superpotential and in the soft supersymmetry breaking scalar potential:

$$\begin{aligned} W &= \mu(\epsilon_i L_i H_2 + H_1 H_2) \\ V_{soft} &= \mu(\epsilon_i B_i L_i H_2 + B H_1 H_2) + m_{L_i H_1}^2 L_i H_1^\dagger + h.c. \end{aligned} \quad (8)$$

Here we have used the same notation for the superfields and their scalar components. A key point to notice is that without the electroweak symmetry breaking, the  $SU(4)$  rotation in the ‘superfields’,  $L_i$  and  $H_1$ ;

$$L_i \rightarrow L_i + \epsilon_i H_1 \quad \text{and} \quad H_1 \rightarrow H_1 - \epsilon_i L_i \quad (9)$$

which gets rid of the  $\epsilon_i$  term (valid up to  $O(\epsilon_i)$ ) leaves invariant the gauge interactions and thus its effect is only to generate the effective couplings as in Eq. (3) with

$$\lambda'_{i33} = \epsilon_i h_b \quad \text{and} \quad \lambda_{i33} = \epsilon_i h_\tau. \quad (10)$$

Under the  $SU(4)$  rotation (9), the scalar potential in (8) becomes

$$V_{soft} = \mu(B H_1 H_2 - \epsilon_i \Delta B_i L_i H_2) + (m_{L_i H_1}^2 - \epsilon_i \Delta m_i^2) L_i H_1^\dagger + h.c. \quad (11)$$

where  $\Delta B_i = B - B_i$  and  $\Delta m_i^2 = m_{H_1}^2 - m_{L_i}^2$ . Eq. (11) shows that the additional lepton number violating mixing mass terms for the ‘scalar fields’  $\tilde{L}_i$  and  $H_{1,2}$  (in the basis of vanishing  $\epsilon_i$ ) arise in the presence of the non-universal soft supersymmetry breaking parameters. Diagonalizing away such mixing mass terms can be made by the following rotation among the scalar fields  $\tilde{L}_i$ ,  $H_1$  and  $H'_2 \equiv i\tau_2 H_2^\dagger$ :

$$\begin{aligned}\tilde{L}_i &\rightarrow \tilde{L}_i - \varepsilon_{i1} H_1 - \varepsilon_{i2} H'_2 \\ H_1 &\rightarrow H_1 + \varepsilon_{i1} \tilde{L}_i \\ H'_2 &\rightarrow H'_2 + \varepsilon_{i2} \tilde{L}_i\end{aligned}\tag{12}$$

where the variables  $\varepsilon_{i1}$  and  $\varepsilon_{i2}$  are determined as

$$\begin{aligned}\varepsilon_{i1} &= \frac{(m_{H_2}^2 + \mu^2 - m_{L_i}^2)(\epsilon_i \Delta m_i^2 - m_{L_i H_1}^2) - \epsilon_i \mu^2 B \Delta B_i}{(m_{H_1}^2 + \mu^2 - m_{L_i}^2)(m_{H_2}^2 + \mu^2 - m_{L_i}^2) - \mu^2 B^2} \\ \varepsilon_{i2} &= \frac{(m_{H_1}^2 + \mu^2 - m_{L_i}^2)\epsilon_i \mu \Delta B_i - \mu B(\epsilon_i \Delta m_i^2 - m_{L_i H_1}^2)}{(m_{H_1}^2 + \mu^2 - m_{L_i}^2)(m_{H_2}^2 + \mu^2 - m_{L_i}^2) - \mu^2 B^2}\end{aligned}\tag{13}$$

As will be discussed later, it is useful to rewrite  $\varepsilon_{i1,2}$  in terms of the variables  $\xi_i$  and  $\eta_i$  defined by

$$\xi_i \equiv \frac{\langle \tilde{\nu}_i \rangle}{\langle H_1 \rangle} - \epsilon_i \quad \text{and} \quad \eta_i \equiv \xi_i + \epsilon_i \frac{\Delta B_i}{B}$$

where  $\langle \tilde{\nu}_i \rangle$  and  $\langle H_1 \rangle$  are the vacuum expectation values of the sneutrino and Higgs boson generated after the electroweak symmetry breaking. Using the minimization condition of the Higgs and sneutrino fields, we obtain

$$\begin{aligned}\varepsilon_{i1} &= -\xi_i - \eta_i \frac{m_A^2 s_\beta^2 (m_{\tilde{\nu}_i}^2 - M_Z^2 c_{2\beta})}{m_{\tilde{\nu}_i}^2 (m_{\tilde{\nu}_i}^2 - m_A^2) - (m_{\tilde{\nu}_i}^2 - m_A^2 s_\beta^2) M_Z^2 c_{2\beta}} \\ \varepsilon_{i2} &= \frac{\eta_i}{t_\beta} \frac{m_A^2 s_\beta^2 m_{\tilde{\nu}_i}^2}{m_{\tilde{\nu}_i}^2 (m_{\tilde{\nu}_i}^2 - m_A^2) - (m_{\tilde{\nu}_i}^2 - m_A^2 s_\beta^2) M_Z^2 c_{2\beta}}\end{aligned}\tag{14}$$

where  $t_\beta = \tan \beta = \langle H_2 \rangle / \langle H_1 \rangle$ . The variables  $\varepsilon_{i1,2}$  control the size of lepton number violating interactions which now arise due to the misalignment between the scalars,  $\tilde{L}_i$  and  $H_{1,2}$ , and fermions,  $L_i$  and  $\tilde{H}_{1,2}$ . That is, the rotation (12) gives rise to the following lepton number violating vertices:

$$\begin{aligned}\mathcal{L}_{eff} &= h_\tau \varepsilon_{i1} \tilde{L}_i L_3 E_3^c + h_b \varepsilon_{i1} \tilde{L}_i Q_3 D_3^c + h_t \varepsilon_{i2} \tilde{L}'_i Q_3 U_3^c \\ &\quad + \frac{g' \varepsilon_{i1}}{\sqrt{2}} [H_1^\dagger L_i \tilde{B} + \tilde{L}_i^\dagger \tilde{H}_1 \tilde{B}] + \frac{g' \varepsilon_{i2}}{\sqrt{2}} \tilde{L}_i^\dagger \tilde{H}_2 \tilde{B} \\ &\quad + \frac{g \varepsilon_{i1}}{\sqrt{2}} [H_1^\dagger \tau^a L_i \lambda^a + \tilde{L}_i^\dagger \tau^a \tilde{H}_1 \lambda^a] + \frac{g \varepsilon_{i2}}{\sqrt{2}} \tilde{L}_i^\dagger \tau^a \tilde{H}_2 \lambda^a + h.c.\end{aligned}\tag{15}$$

where  $\tilde{L}'_i \equiv i\tau_2 \tilde{L}_i^\dagger$ ,  $\tau^a$  are Pauli matrices and  $\lambda^a$  represent the  $SU(2)$  gauginos. Applying the constraint (7) to the couplings in Eqs. (10) and (15), we get [25]

$$\begin{aligned}\epsilon_i &< 1.2 \times 10^{-5} c_\beta \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^{1/2} \\ \varepsilon_{i1} &< 3 \times 10^{-7} \left( \frac{m_{\chi^0}}{300 \text{ GeV}} \right)^{1/2} \\ \varepsilon_{i2} &< 2 \times 10^{-7} s_\beta \left( \frac{m_{L_i}}{300 \text{ GeV}} \right)^{1/2}\end{aligned}\tag{16}$$

where  $\tilde{m}$  is the smallest mass of the sfermions involved in the  $\lambda'_{i33}$  term;  $L_i, Q_3$  and  $D_3^c$ ,  $m_\chi$  is a gaugino mass involved in the process  $\chi \rightarrow L_i H_1$  and the last equation comes from the process  $\tilde{L}_i \rightarrow Q_3 U_3^c$ .

The sizes of certain bilinear parameters are determined to generate realistic neutrino masses and mixing in our bilinear model. First of all, upon electroweak symmetry breaking, the Higgs and sneutrino acquire vacuum expectation values and generate a tree-level neutrino mass matrix

$$M_{ij}^{tree} = \frac{M_Z^2}{F_N} \xi_i \xi_j c_\beta^2 \tag{17}$$

where  $F_N = M_1 M_2 / (c_W^2 M_1 + s_W^2 M_2) + M_Z^2 c_{2\beta} / \mu$  [13]. Recall that  $\xi_i$  arises through the mismatch of soft terms between  $L_i$  and  $H_1$  as follows;

$$\xi_i = \epsilon_i \frac{\Delta m_i^2 + \Delta B_i \mu t_\beta - m_{L_i H_1}^2 / \epsilon_i}{m_{\nu_i}^2}. \tag{18}$$

The tree mass in Eq. (17) gives the heavier mass scale,  $m_3 = \frac{M_Z^2}{F_N} \xi^2 c_\beta^2$ . Considering the atmospheric neutrino mass-squared difference,  $\Delta m_{\text{atm}}^2 \approx 2.5 \times 10^{-3} \text{ eV} \approx m_3^2$ , we get

$$\xi c_\beta = 7.4 \times 10^{-7} \left( \frac{F_N}{M_Z} \right)^{1/2} \left( \frac{m_3}{0.05 \text{ eV}} \right)^{1/2} \tag{19}$$

Since the two mixing angles,  $\theta_{23} = \theta_{\text{atm}}$  and  $\theta_{13}$ , satisfy

$$\tan \theta_{23} = \xi_2 / \xi_3 \approx 1, \quad |\tan \theta_{13}| = |\xi_1| / \sqrt{\xi_2^2 + \xi_3^2} \ll 1 \tag{20}$$

we need  $\xi_1 < 0.3 \xi_{2,3}$  to make small  $\theta_{13}$  and  $\xi_2 \approx \xi_3$  for near maximal atmospheric mixing.

Thus, current neutrino oscillation data require

$$\xi_1 \ll \xi_2 \approx \xi_3 \approx 5.2 \times 10^{-7} \frac{1}{c_\beta} \left( \frac{F_N}{M_Z} \right)^{1/2} \left( \frac{m_3}{0.05 \text{ eV}} \right)^{1/2} \tag{21}$$

Let us now consider how one-loop corrections generate the neutrino masses and mixing accounting for the solar neutrino oscillation. In the bilinear model, the bi-large mixing of the atmospheric and solar neutrinos cannot be obtained under the assumption of universal soft terms [13]. Thus, one needs to introduce non-universality in soft terms in order to accommodate the large solar mixing.

Depending on the degrees of the deviation from the universality, we can consider two cases. First, the non-universality of soft parameters can arise due to small mismatches (likely to be caused by some threshold corrections) in the renormalization group evolution. In this case, the quantities  $\Delta m_i^2$ ,  $m_{L_i H_1}^2/\epsilon_i$  and  $\mu\Delta B_i$  are much smaller than the typical soft mass-squared  $\tilde{m}^2$  so that the induced trilinear couplings in Eq. (10) give the major contribution to the size of  $m_2 \approx \sqrt{\Delta m_{\text{sol}}^2}$  [14, 16]. As discussed before, this causes the contradiction of Eqs. (5) and (7). In other words, the condition of  $m_2 \sim 8$  meV yields  $\epsilon_i \sim 4 \times 10^{-3} c_\beta$  which is far above the first constraint in Eq. (16).

However, we point out that there is a different way of reconciling the neutrino data with the baryogenesis requirement. Note that one cannot exclude the possibility that the loop mass is larger than the tree mass. For instance, one can take the superpotential bilinear parameter  $\epsilon_i$  much larger than  $\xi_i$ , accepting a very small deviation of the non-universality or a cancellation among the terms [5], see Eq. (18). In this situation, the heavier neutrino mass scale can be produced mainly by the bottom-sbottom loop which can be rewritten from Eqs. (4) and (10) as follows:

$$M_{ij}^{\text{loop}} = \frac{3h_b^2}{8\pi^2} \epsilon_i \epsilon_j \frac{m_b^2 (A_b + \mu \tan \beta)}{m_{\tilde{b}_1}^2 - m_{\tilde{b}_2}^2} \ln \frac{m_{\tilde{b}_1}^2}{m_{\tilde{b}_2}^2}. \quad (22)$$

As the above loop contribution determines the atmospheric neutrino mass and mixing, the condition (21) has to be replaced by

$$\epsilon_1 \ll \epsilon_2 \approx \epsilon_3 \approx 8 \times 10^{-3} c_\beta \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^{1/2}. \quad (23)$$

Similarly to the previous discussions,  $\epsilon_{2,3}$  cannot satisfy the baryogenesis constraint (16) at all. But,  $\epsilon_1$  can be made arbitrarily small. Let us recall that it is sufficient to suppress lepton number violating couplings for just one lepton flavour. In our case, it is the electron number, which is implied by the smallness of  $\theta_{13}$ . Now, in order for the tree mass (17) to produce the solar neutrino mass and mixing, we need

$$\xi_1 \sim \xi_2 \sim 3 \times 10^{-7} \frac{1}{c_\beta} \left( \frac{F_N}{M_Z} \right)^{1/2}. \quad (24)$$

For such small  $\epsilon_1$  and the small deviation of the universality, we expect  $\xi_1 \simeq \eta_1$  and thus the variables  $\varepsilon_{11,12}$  in Eq. (14) can be approximated by

$$\varepsilon_{11} \approx -\xi_1 \frac{m_{\tilde{\nu}_1}^2 - m_A^2 c_\beta^2}{m_{\tilde{\nu}_1}^2 - m_A^2}, \quad \varepsilon_{12} \approx -\frac{\xi_1}{t_\beta} \frac{m_A^2 s_\beta^2}{m_{\tilde{\nu}_1}^2 - m_A^2}, \quad (25)$$

neglecting  $M_Z^2$  terms. From this, one can see that the baryogenesis constraint (16) can be satisfied if  $1 < t_\beta < (m_{\tilde{\nu}_1}^2 - m_A^2)/(m_{\tilde{\nu}_1}^2 - m_A^2 c_\beta^2)$ .

Secondly, we consider the more general non-universality implying that  $\Delta m_i^2$ ,  $m_{L_i H_1}^2/\epsilon_i$  and  $\mu \Delta B_i$  are of the order  $\tilde{m}^2$ . In this case, the neutral scalar and neutralino exchange loops can give important contributions to the one-loop mass as long as  $\tan \beta$  is not too large and the large misalignment between  $\xi_i$  and  $\eta_i$  is allowed. Adopting the result of Ref. [15], the one-loop mass coming from the neutral scalar loops is roughly given by

$$M_{ij}^{loop} \approx \frac{g^2}{64\pi^2} m_{\chi^0} \theta_{i\phi} \theta_{j\phi} B_0(m_{\chi^0}^2, m_\phi^2) \quad (26)$$

where  $B_0(x, y) = -\frac{x}{x-y} \ln \frac{x}{y} - \ln \frac{x}{Q^2} + 1$  and  $\phi$  represents the neural Higgs bosons,  $h, H$  and  $A$ . Neglecting unimportant contribution of  $\xi_i$ , the variables  $\theta_{i\phi}$  are approximately given by

$$\begin{aligned} \theta_{ih} &\approx \eta_i s_\beta m_A^2 \frac{m_{\tilde{\nu}_i}^2 c_{\alpha-\beta} - M_Z^2 c_{2\beta} c_{\alpha+\beta}}{(m_{\tilde{\nu}_i}^2 - m_h^2)(m_{\tilde{\nu}_i}^2 - m_H^2)} \\ \theta_{iH} &\approx \eta_i s_\beta m_A^2 \frac{m_{\tilde{\nu}_i}^2 s_{\alpha-\beta} - M_Z^2 c_{2\beta} s_{\alpha+\beta}}{(m_{\tilde{\nu}_i}^2 - m_h^2)(m_{\tilde{\nu}_i}^2 - m_H^2)} \\ \theta_{iA} &\approx i\eta_i s_\beta \frac{m_A^2}{m_A^2 - m_{\tilde{\nu}_i}^2} \end{aligned} \quad (27)$$

where  $m_{h,H}$  are the Higgs boson masses at tree-level determined by  $m_{h,H}^2 = 1/2[m_A^2 + M_Z^2 \mp \sqrt{(m_A^2 + M_Z^2)^2 - 4m_A^2 M_Z^2 c_{2\beta}^2}]$ , and the angle  $\alpha$  is defined by  $c_{2\alpha} = c_{2\beta}(m_A^2 - M_Z^2)/(m_h^2 - m_H^2)$  and  $s_{2\alpha} = s_{2\beta}(m_A^2 + M_Z^2)/(m_h^2 - m_H^2)$ . Our convention for the pseudo-scalar Higgs boson mass is that  $m_A^2 = -\mu B/c_\beta s_\beta$ . Requiring  $m_2 \sim M_{ij}^{loop}$ , one obtains

$$\theta_{i\phi} \sim 6 \times 10^{-6} \left( \frac{300 \text{ GeV}}{m_{\chi^0}} \right)^{1/2} \left( \frac{m_\phi}{m_{\chi^0}} \right) \left( \frac{m_2}{8 \text{ meV}} \right) \quad (28)$$

As discussed in Ref. [15], the large mixing of solar neutrinos require  $\theta_{1\phi} \sim \theta_{2\phi}$ . This has to be contrasted the condition  $\xi_2 \approx \xi_3$  (21) for the large atmospheric neutrino mixing.

From Eqs. (16), (21) and (28), one sees that the couplings  $\varepsilon_{i1,i2}$  are required to be smaller than  $\xi_i$  or  $\theta_{i\phi}$  by one-order of magnitude. Thus, it is generally difficult to satisfy both the

baryogenesis constraints and obtain the realistic neutrino masses and mixing. However, it is not impossible to find some reasonable parameter space where both requirements are reconciled, which is due to the fact that the variables  $\varepsilon_{i1,i2}$  and  $\xi_i$  or  $\theta_{i\phi}$  have different dependencies on the input parameters. Comparing Eq. (14) with Eq. (27), we notice that  $\varepsilon_{i1,i2}$  (or  $\xi_i$  and  $\eta_i$ ) can be made small while keeping  $\theta_{i\phi} \sim 6 \times 10^{-6}$  (28) when the sneutrino mass  $m_{\tilde{\nu}_i}$  is close to one of the Higgs boson masses,  $m_h, m_H$  and  $m_A$ . Since the heavy Higgs scalar mass,  $m_H$ , is usually close to the pseudo scalar mass,  $m_A$ , and  $s_\beta^2 \approx 1$ , we find it better that the sneutrino mass is closer to the light Higgs scalar mass, that is,  $m_{\tilde{\nu}_1} \sim m_h$ . Barring cancellation, both terms in  $\varepsilon_{i1}$  (14) should be less than  $3 \times 10^{-7}$ . Again here, this is possible for  $i = 1$ , that is, the electron number violating parameters,  $\varepsilon_{11,12}$ , can only be suppressed for our purpose. For illustration, let us calculate  $\varepsilon_{i1} + \xi_i$ ,  $\varepsilon_{i2}$  and  $\theta_{i\phi}$  for the cases with  $m_A = 100, 300$  GeV and  $\tan \beta = 3, 30$ . In what follows, we present the values of  $(\theta_{ih}, \theta_{iH}, \theta_{iA}; \varepsilon_{i1} + \xi_i, \varepsilon_{i2})$  normalized with  $\varepsilon_{i2} = 1$ , indicating the rough ranges of  $m_{\tilde{\nu}_i}$  allowing for  $\varepsilon_{i1} > \theta_{i\phi}/20$ :

$$\text{Case 1} \quad t_\beta = 3, m_A = 100 \text{ GeV}, m_h = 60 \text{ GeV} \quad (29)$$

$$(+77, +64, -9.1; -3.6, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 55 \text{ GeV}$$

$$(-22, -21, -6.5; -0.96, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 71 \text{ GeV}$$

$$\text{Case 2} \quad t_\beta = 30, m_A = 100 \text{ GeV}, m_h = 90 \text{ GeV} \quad (30)$$

$$(-113, +342, -77; -17, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 73 \text{ GeV}$$

$$(-66, -210, -49; +11, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 115 \text{ GeV}$$

$$\text{Case 3} \quad t_\beta = 3, m_A = 300 \text{ GeV}, m_h = 72 \text{ GeV} \quad (31)$$

$$(+29, +53, -8.4; -1.5, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 60 \text{ GeV}$$

$$(-10, -26, -5.5; +0.54, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 90 \text{ GeV}$$

$$\text{Case 4} \quad t_\beta = 30, m_A = 300 \text{ GeV}, m_h = 91 \text{ GeV} \quad (32)$$

$$(+18, +289, -74; -14, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 75 \text{ GeV}$$

$$(-6.9, -212, -49; +11, 1) \quad \text{for} \quad m_{\tilde{\nu}_i} = 115 \text{ GeV}$$

From the above calculation, one can see that the non-erasure condition can be satisfied if the difference between the sneutrino and the light Higgs boson mass is within 10%. In order to confirm the above properties, we made a numerical calculation to find a set of points satisfying both the baryogenesis constraints and the atmospheric and solar neutrino data.

For this, we incorporate the exact formulae for the neutrino mass matrix derived in Ref. [15]. In Figures 1 and 2, we plot the variable  $\varepsilon_{11}$  in terms of the electron sneutrino mass  $m_{\tilde{\nu}_1}$  for all the points accommodating all the observed neutrino data for Cases 1 and 2. The plots clearly show the suppression of  $\varepsilon_{11}$  when the sneutrino mass is close to a Higgs boson mass. Similar behavior is also found in Cases 3 and 4.

Another way of suppressing  $\varepsilon_{i1}$  is to arrange a cancellation between two terms in  $\varepsilon_{i1}$ . From Eqs. (14) and (27), one generally has

$$\varepsilon_{i1} \sim -\xi_i - t_\beta \varepsilon_{i2} \quad \text{and} \quad \theta_{i\phi} \sim t_\beta \varepsilon_{i2} \quad (33)$$

for  $m_{\tilde{\nu}_i} \gg M_Z$ . Now, one can see that the conditions (16) and (28) can be satisfied for  $t_\beta \sim 30$  with the cancellation in  $\varepsilon_{i1} \sim \xi_i + \theta_{i\phi}$ . Again, this can work only for the electron direction with  $\xi_1 \sim \theta_{1\phi}$  since Eq. (21) shows  $\xi_i \gg \theta_{i\phi} \sim 6 \times 10^{-6}$  for  $i = 2, 3$  and large  $\tan \beta$ .

In conclusion, we have investigated how the cosmological requirement for a successful baryogenesis can be reconciled with a realistic neutrino mass matrix in the R-parity violating version of supersymmetric standard model. Our main focus has been to see whether the  $B - L$  violating interactions can be sufficiently suppressed in order not to erase a pre-existing baryon or lepton asymmetry of the universe. Such a baryogenesis constraint cannot be satisfied if the trilinear R-parity violating couplings are introduced to explain the atmospheric and solar neutrino masses and mixing under the assumption of the universal soft supersymmetry breaking masses. In the bilinear model, the observed neutrino data can be well explained if the non-universality is allowed. Our analysis shows that the non-erasure condition can be met by suppressing the electron number violating parameters, which is related to the smallness of the angle  $\theta_{13}$ . In the case of a large violation of the universality, the electron sneutrino mass has to be nearly degenerate with the light Higgs scalar mass. For a small violation of the universality, we argued that the situation of the loop mass dominating over the tree mass is preferred contrary to the usual consideration. A consequence of our analysis is that the bilinear R-parity violating supersymmetric standard model can provide a framework not only for a realistic neutrino mass matrix but also for a successful baryogenesis through the Affleck-Dine Mechanism. Finally, let us note that our consideration is not relevant if the electro-weak baryogenesis is operative.

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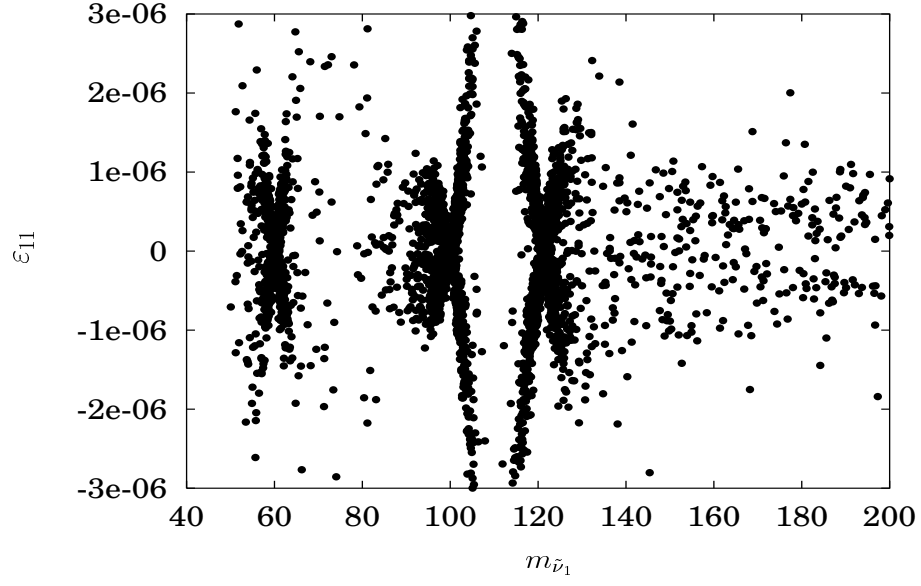


FIG. 1: The quantity  $\varepsilon_{11}$  is shown as a function of the electron sneutrino mass  $m_{\tilde{\nu}_1}$  for all the points generating the required neutrino masses and mixing for Case 1 in Eq. (29).

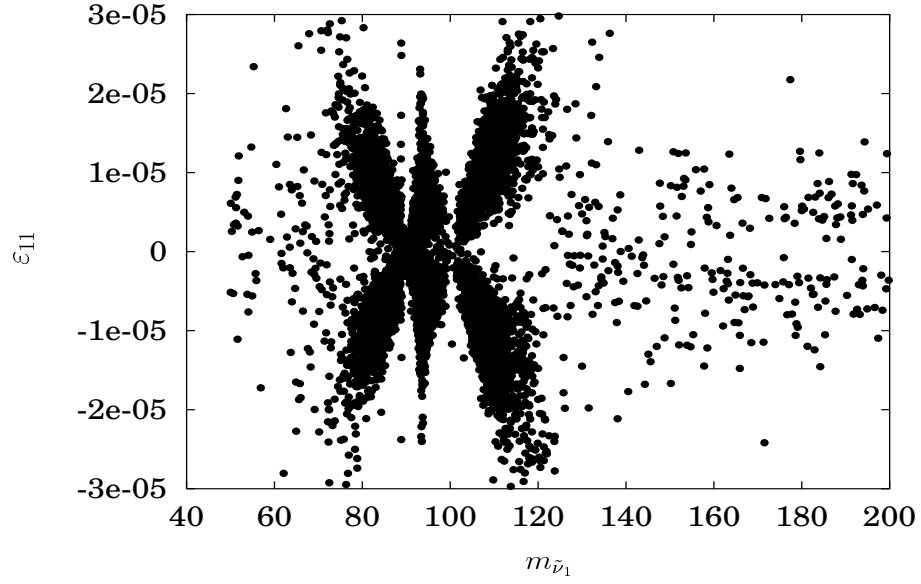


FIG. 2: Same as FIG. 1 for Case 2 in Eq. (30).